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Quarkonium bound state equation in the Wilson approach with minimal surfaces[1]

F. Jugeau[a] and H. Sazdjian[b]

[a] *Laboratoire de Physique Théorique*[2],
Université Paris XI, Bât. 210, F-91405 Orsay Cedex, France
E-mail: Frederic.Jugeau@th.u-psud.fr

[b] *Groupe de Physique Théorique, Institut de Physique Nucléaire*[3],
Université Paris XI, F-91406 Orsay Cedex, France
E-mail: sazdjian@ipno.in2p3.fr

Abstract

Wilson loop averages are evaluated for large contours and in the large-N_c limit by means of minimal surfaces. This allows the study of the quark-antiquark gauge invariant Green function through its dependence on Wilson loops. A covariant bound state equation is derived which in the center-of-mass frame and at equal-times takes the form of a Breit–Salpeter type equation. The interaction potentials reduce in the static case to a confining linear vector potential. For moving quarks, flux tube like contributions are present. The nonrelativistic limit is considered.

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[2] Unité Mixte de Recherche 8627.

[3] Unité Mixte de Recherche 8608.

The Wilson loop approach [1] is based on the use of the gluon field path-ordered phase factor along a line C joining a point x to a point y:

$$U(C_{yx}, y, x) \equiv U(y, x) = Pe^{-ig\int_x^y dz^\mu A_\mu(z)}. \tag{1}$$

Equations satisfied by phase factors were obtained and analyzed by Mandelstam [2] and Nambu [3].

The Wilson loop is defined as the trace in color space of the phase factor on a closed contour C:

$$\Phi(C) = \frac{1}{N_c}\mathrm{tr}Pe^{-ig\oint_C dx^\mu A_\mu(x)}. \tag{2}$$

Its vacuum expectation value is denoted $W(C)$:

$$W(C) = \langle\Phi(C)\rangle. \tag{3}$$

Loop equations were obtained by Polyakov [4] and Makeenko and Migdal [5, 6, 7, 8]. The Wilson loop essentially satisfies two types of equation, which are equivalent to the QCD equations of motion:

1) The Bianchi identity.

2) The loop equations (or Makeenko–Migdal equations). Those actually represent an infinite chain of coupled equations.

A third property, factorization, is obtained in the large-N_c limit for two disjoint contours:

$$W(C_1, C_2) = W(C_1)W(C_2). \tag{4}$$

Considerable simplification occurs in the large-N_c limit of the theory [9]. In that limit, for large contours, i.e., at large distances, nonperturbative asymptotic solutions to the Wilson loops are represented by the minimal surfaces having as supports the loop contours [5, 6]. Among various types of surfaces, minimal surfaces are the only ones that satisfy the Bianchi identity; actually the latter becomes identical to the defining equation of minimal surfaces [10]. Furthermore, if at large distances one ignores short-distance perturbative solutions, then the unrenormalized coupling constant which appears in the loop equations can be adjusted in relation with the string tension [10] and therefore minimal surfaces define on their own independent solutions. That property fails, however, when short-distance perturbative solutions are taken into account, since in this case the unrenormalized coupling constant is adjusted with respect to the perturbative regime. Minimal surfaces then become only large-distance asymptotic solutions. Nevertheless, if one is interested only in the large-distance behavior of the theory, saturation of the Wilson loop averages by minimal surfaces provides a correct description of the theory in

this regime. In that case, the Wilson loop average can be represented by the following functional of the contour C:

$$W(C) = e^{-i\sigma A(C)}, \tag{5}$$

where σ is the string tension and $A(C)$ the minimal area with contour C.

Minimal surfaces also appear as natural solutions to the Wilson loop averages in two-dimensional gauge theories [11].

To deal with the quarkonium bound state problem, we start with the two-particle gauge invariant Green function for quarks q_1 and q_2 with different flavors and with masses m_1 and m_2:

$$G(x_1, x_2; x_1', x_2') \equiv \langle \overline{\psi}_2(x_2) U(x_2, x_1) \psi_1(x_1) \overline{\psi}_1(x_1') U(x_1', x_2') \psi_2(x_2') \rangle_{A, q_1, q_2}. \tag{6}$$

Here $U(x_2, x_1)$ is the phase factor (1) taken along the straight-line $x_1 x_2$ (and similarly for $U(x_1', x_2')$). Integrating in the large-N_c limit with respect to the quark fields, one obtains:

$$G(x_1, x_2; x_1', x_2') = -\langle \mathrm{tr}\, U(x_2, x_1) S_1(x_1, x_1') U(x_1', x_2') S_2(x_2', x_2) \rangle_A, \tag{7}$$

where S_1 and S_2 are the quark and antiquark propagators in the presence of the external gluon field and tr designates the trace with respect to the color group.

The Green function G satisfies the following equation with respect to the Dirac operator of particle 1 acting on x_1:

$$\begin{aligned}(i\gamma.\partial_{(x_1)} - m_1) G(x_1, x_2; x_1', x_2') &= -i \langle \mathrm{tr}\, U(x_2, x_1) \delta^4(x_1 - x_1') U(x_1', x_2') S_2(x_2', x_2) \rangle_A \\ -i\gamma^\alpha \langle \mathrm{tr} \int_0^1 d\sigma (1-\sigma) &\frac{\delta U(x_2, x_1)}{\delta x^\alpha(\sigma)} S_1(x_1, x_1') U(x_1', x_2') S_2(x_2', x_2) \rangle_A,\end{aligned} \tag{8}$$

where the segment $x_1 x_2$ has been parametrized with the parameter σ as $x(\sigma) = (1 - \sigma) x_1 + \sigma x_2$; furthermore, the operator $\delta / \delta x^\alpha$ does not act on the explicit boundary point x_1 of the segment, this contribution having been cancelled by the contribution of the gluon field A coming from the quark propagator S_1. A similar equation also holds with the Dirac operator of particle 2 acting on x_2. The operation $\delta U(x_2, x_1) / \delta x(\sigma)$ introduces an insertion of the field strength F at the point $x(\sigma)$ of the straight-line $x_1 x_2$ [2, 4, 8].

In order to make apparent the Wilson loop structure of the two-particle Green function, we adopt a representation for the quark propagator in the external gluon field based on an explicit use of the phase factor along straight lines [10]. Introducing the gauge covariant composite object $\widetilde{S}(x, x')$, made of a free fermion propagator $S_0(x - x')$ (without color group content) multiplied by the path-ordered phase factor $U(x, x')$ [Eq. (1)] taken along the straight segment $x' x$,

$$\widetilde{S}(x, x') \equiv S_0(x - x') U(x, x'), \tag{9}$$

one shows that the quark propagator $S(x,x')$ in the external gluon field satisfies the following functional integral equation in terms of $\widetilde{S}$:

$$S(x,x') = \widetilde{S}(x,x') - \int d^4x'' S(x,x'')\gamma^\alpha \int_0^1 d\lambda\, \lambda \frac{\delta}{\delta x^\alpha(\lambda)} \widetilde{S}(x'',x'), \tag{10}$$

where the operator $\delta/\delta x^\alpha(\lambda)$ acts on the factor U of $\widetilde{S}$, along the internal part of the segment $x'x''$, with x' held fixed. A similar equation in which the roles of x and x' are interchanged also holds. Those equations lead to iteration series for S in which the gauge covariance property is maintained at each order of the iteration.

Using the above representations for the quark propagators in Eq. (7) one obtains for the two-particle Green function a series expansion where each term contains a Wilson loop along a skew-polygon:

$$G = \sum_{i,j=1}^{\infty} G_{i,j}, \tag{11}$$

where $G_{i,j}$ represents the contribution of the term of the series having $(i-1)$ points of integration between x_1 and x_1' (i segments) and $(j-1)$ points of integration between x_2 and x_2' (j segments). We designate by $C_{i,j}$ the contour associated with the term $G_{i,j}$. A typical configuration for the contour of $G_{4,3}$ is represented in Fig. 1.

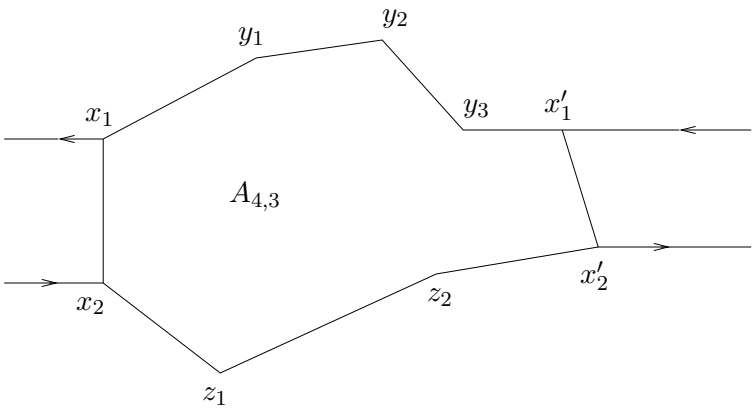

Figure 1: Contour $C_{4.3}$ associated with the term $G_{4,3}$. $A_{4,3}$ is the minimal surface with contour $C_{4,3}$.

Each segment of the quark lines supports a free quark propagator and except for the first segments (or the last ones, depending on the representation that is used) the Wilson loop is submitted to one functional derivative on each such segment. One then uses for the averages of the Wilson loops appearing in the above series the representation with minimal surfaces [Eq. (5)].

Representation (10) for the quark propagator is also used in Eq. (8) and its partner satisfied by the two-particle Green function G. One obtains two compatible equations for G where the right-hand sides involve the series of the terms $G_{i,j}$ of Eq. (11) and their functional derivative along the segment x_1x_2. In order to obtain bound state equations, it is necessary to reconstruct in the right-hand sides the bound state poles contained in G [12]. In x-space, bound states are reached by taking the large separation time limit between the pair of points (x_1, x_2) and (x'_1, x'_2) [13].

To produce a bound state pole, it is necessary that there be a coherent sum of contributions coming from each $G_{i,j}$, since the latter, taken individually, do not have poles. It is evident from representation (5) that functional derivatives acting on Wilson loop averages give rise to functional derivatives of the corresponding minimal surfaces $A_{i,j}$. One then separates the various contributions of the right-hand sides of Eq. (8) and of its partner into different categories having the property of irreducibility. The first category corresponds to terms in which the functional derivative along the segment x_1x_2 acts alone on a given A. In the second category, the derivative along x_1x_2 is accompanied by another derivative along one of the segments of the quark lines, the two acting on the same A. In the third category, the derivative along x_1x_2 and two other derivatives along segments of the quark lines act on the same A or on the product of two As, respecting the irreducibility property of kernels, in the sense that they are not parts of the series expansion of a factorized G, and so forth.

Let us now consider the terms of the first category. One notices that the derivative along the segment x_1x_2 acts on areas $A_{i,j}$ with contour $C_{i,j}$ which are different from one term of the series to the other (the number of segments being different). To have a coherent sum of those contributions it is necessary to expand each such derivative around the derivative of the lowest-order contour $C_{1,1}$, represented in Fig. 2.

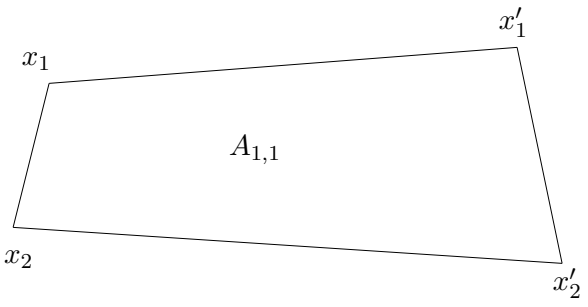

Figure 2: The lowest-order contour $C_{1,1}$ and its minimal surface $A_{1,1}$.

It is that term that can be factorized and can lead through the summation of the factored series to the reappearance of the Green function G and to its poles. The remaining

terms do not lead to pole terms. Similarly, in the second category of terms, the two derivative contributions containing the derivative along x_1x_2 should be expanded around the lowest-order contribution coming from the contours $C_{2,1}$ or $C_{1,2}$, and so forth.

In general, the derivative of the areas along x_1x_2 depends among others on the slope of the areas in the orthogonal direction to x_1x_2, thus feeling the positions of the other points x_1' and x_2'. In order to obtain bound state equations depending solely on x_1 and x_2 and not on the remote points x_1' and x_2', it is necessary to associate the slopes of the areas along the segment x_1x_2 with the quark momenta. Taking then the large separation time limit, one ends up with two covariant compatible bound state equations, in which the interaction kernels or potentials are given by various functional derivatives involving at least one derivative along the segment x_1x_2.

The fact that one has two independent but compatible equations means that the relative time variable does not play here a dynamical role and could in principle be eliminated. That is done by taking the difference and sum of the two equations. One finally ends up with one single three-dimensional equation at equal-times in the center-of-mass frame, having the structure of a Breit–Salpeter type equation [14, 15]. Keeping for the potentials the terms containing one functional derivative of the area $A_{1,1}$ [Fig. 2], the equation takes the form [10]

$$\Big[P_0 - (h_{10} + h_{20}) - \gamma_{10}\gamma_1^\mu A_{1\mu} - \gamma_{20}\gamma_2^\mu A_{2\mu}\Big]\psi(\mathbf{x}) = 0, \tag{12}$$

where ψ is a 4×4 matrix wave function of the relative coordinate $x = x_2 - x_1$ considered at equal times, P_0 the center-of-mass total energy and h_{10} and h_{20} the quark and antiquark Dirac hamiltonians; the Dirac matrices of the quark (with index 1) act on ψ from the left, while the Dirac matrices of the antiquark (with index 2) act on ψ from the right. The potentials A_1 and A_2 are defined through the equations

$$A_{1\mu} = \sigma \int_0^1 d\sigma'(1-\sigma')\frac{\delta A_{1,1}}{\delta x^\mu(\sigma')}, \qquad A_{2\mu} = \sigma \int_0^1 d\sigma'\,\sigma'\frac{\delta A_{1,1}}{\delta x^\mu(\sigma')}, \tag{13}$$

$x(\sigma')$ belonging to the segment x_1x_2.

The time components of A_1 and A_2 add up in the wave equation. For their sum, one has the expression (in the c.m. frame)

$$\begin{aligned} A_{10} + A_{20} = \sigma r \frac{E_1E_2}{E_1+E_2}\Big\{&\Big(\frac{E_1}{E_1+E_2}\epsilon(p_{10}) + \frac{E_2}{E_1+E_2}\epsilon(p_{20})\Big) \\ &\times\sqrt{\frac{r^2}{\mathbf{L}^2}}\left(\arcsin\Big(\frac{1}{E_2}\sqrt{\frac{\mathbf{L}^2}{r^2}}\Big) + \arcsin\Big(\frac{1}{E_1}\sqrt{\frac{\mathbf{L}^2}{r^2}}\Big)\right) \\ +(\epsilon(p_{10}) - \epsilon(p_{20}))&\Big(\frac{E_1E_2}{E_1+E_2}\Big)\Big(\frac{r^2}{\mathbf{L}^2}\Big)\left(\sqrt{1-\frac{\mathbf{L}^2}{r^2E_2^2}} - \sqrt{1-\frac{\mathbf{L}^2}{r^2E_1^2}}\right)\Big\}. \end{aligned} \tag{14}$$

Here, $r=\sqrt{\mathbf{x}^2}$, $E_a=\sqrt{m_a^2+\mathbf{p}^2}$, $a=1,2$, with m_a the quark masses, $\mathbf{p}$ the c.m. momentum, $\mathbf{p}=(\mathbf{p}_2-\mathbf{p}_1)/2$, $\mathbf{L}$ the c.m. orbital angular momentum, and $\epsilon(p_{10})$ and $\epsilon(p_{20})$ the energy sign operators of the free quark and the antiquark, respectively:

$$\epsilon(p_{a0})=\frac{h_{a0}}{E_a},\qquad a=1,2. \tag{15}$$

The space components of A_1 and A_2 are orthogonal to $\mathbf{x}$. The expression of $\mathbf{A}_1$ is (in the c.m. frame):

$$\begin{aligned}\mathbf{A}_1 \;=\; & -\sigma r\frac{E_1E_2}{E_1+E_2}\Big\{\frac{r^2}{2\mathbf{L}^2}\frac{E_1E_2}{E_1+E_2}\mathbf{p}^t \\ & \times\sqrt{\frac{r^2}{\mathbf{L}^2}}\left(\arcsin\Big(\frac{1}{E_2}\sqrt{\frac{\mathbf{L}^2}{r^2}}\Big)+\arcsin\Big(\frac{1}{E_1}\sqrt{\frac{\mathbf{L}^2}{r^2}}\Big)\right) \\ & +\frac{1}{E_2}\mathbf{p}^t\Big(\frac{E_1E_2}{E_1+E_2}\Big)\Big(\frac{r^2}{\mathbf{L}^2}\Big)\left(\sqrt{1-\frac{\mathbf{L}^2}{r^2E_2^2}}-\sqrt{1-\frac{\mathbf{L}^2}{r^2E_1^2}}\right) \\ & -\frac{1}{2}\mathbf{p}^t\Big(\frac{r^2}{\mathbf{L}^2}\Big)\left(\frac{E_1}{E_1+E_2}\sqrt{1-\frac{\mathbf{L}^2}{r^2E_2^2}}+\frac{E_2}{E_1+E_2}\sqrt{1-\frac{\mathbf{L}^2}{r^2E_1^2}}\right)\Big\}.\end{aligned} \tag{16}$$

Here, $\mathbf{p}^t$ is the transverse part of $\mathbf{p}$ with respect to $\mathbf{x}$:

$$\mathbf{p}^t=\mathbf{p}-\mathbf{x}\frac{1}{\mathbf{x}^2}\mathbf{x}.\mathbf{p}. \tag{17}$$

The expression of $\mathbf{A}_2$ is obtained from that of $\mathbf{A}_1$ by an interchange in the latter of the indices 1 and 2 and a change of sign of $\mathbf{p}^t$.

For sectors of quantum numbers where $\mathbf{L}^2=0$, the expressions of the potentials become:

$$A_{10}+A_{20}=\frac{1}{2}(\epsilon(p_{10})+\epsilon(p_{20}))\sigma r, \tag{18}$$

$$\mathbf{A}_1=-\frac{1}{E_1E_2}\Big(\frac{1}{3}(E_1+E_2)-\frac{1}{2}E_1\Big)\mathbf{p}^t\sigma r,$$

$$\mathbf{A}_2=+\frac{1}{E_1E_2}\Big(\frac{1}{3}(E_1+E_2)-\frac{1}{2}E_2\Big)\mathbf{p}^t\sigma r. \tag{19}$$

The potentials are generally momentum dependent operators and necessitate an appropriate ordering of terms.

From the structure of the wave equation (12) and the expressions of the potentials, one deduces that the interaction is confining and of the vector type. However, compared to the conventional timelike vector potential, it has additional pieces of terms contributing to the orbital angular momentum dependent parts. A closer analysis of those terms shows that they can be interpreted as being originated from the moments of inertia of the segment x_1x_2 carrying a constant linear energy density equal to the string tension. The interaction

potentials are therefore provided by the energy-momentum vector of the segment joining the quark to the antiquark, in similarity with the color flux tube picture of confinement. An analogous equation had also been proposed by Olsson *et al.* on the basis of a model where the quarks are attached at the ends of a straight string or a color flux tube [16, 17]. A similar conclusion had also been reached by Brambilla, Prosperi *et al.* on the basis of the analysis of the relativistic corrections to the nonrelativistic limit of the Wilson loop [18, 19, 20].

For heavy quarks, one can expand equation (12) around the nonrelativistic limit. To order $1/c^2$, the hamiltonian becomes (in the c.m. frame):

$$\begin{aligned} H &= \frac{\mathbf{p}^2}{2\mu} + \sigma r - \frac{2\hbar\sigma}{\pi}\Big(\frac{1}{m_1}+\frac{1}{m_2}\Big) - \frac{1}{8}\Big(\frac{1}{m_1^3}+\frac{1}{m_2^3}\Big)(\mathbf{p}^2)^2 + \frac{\hbar^2}{4}\Big(\frac{1}{m_1^2}+\frac{1}{m_2^2}\Big)\frac{\sigma}{r} \\ &\quad -\frac{\sigma}{6r}\Big(\frac{1}{m_1^2}+\frac{1}{m_2^2}-\frac{1}{m_1 m_2}\Big)(\mathbf{L}^2+2\hbar^2) + \frac{\sigma}{2r}\Big(\frac{\mathbf{L}.\mathbf{s}_1}{m_1^2}+\frac{\mathbf{L}.\mathbf{s}_2}{m_2^2}\Big) \\ &\quad -\frac{2\sigma}{3r}\Big(\frac{1}{m_1^2}-\frac{1}{2m_1 m_2}\Big)\mathbf{L}.\mathbf{s}_1 - \frac{2\sigma}{3r}\Big(\frac{1}{m_2^2}-\frac{1}{2m_1 m_2}\Big)\mathbf{L}.\mathbf{s}_2. \end{aligned} \tag{20}$$

[$\mu = m_1 m_2/(m_1+m_2)$, $\mathbf{s}_1$ and $\mathbf{s}_2$ are the spin operators of the quark and of the antiquark.] Several remarks can be made at this stage. First, the hamiltonian is independent of spin-spin interactions. Second, purely orbital angular momentum dependent pieces (proportional to $\mathbf{L}^2$) are present, the origin of which is related to the contribution to the rotational motion of the system of the moments of inertia of the flux tube, represented by the straight segment joining the quark to the antiquark. Those terms were also obtained in Refs. [18] and [16, 17]. Third, two kinds of spin-orbit term are present. The first comes from the contribution of a conventional timelike vector interaction represented by the potential σr, which is the dominant part of the combination $A_{10}+A_{20}$ [Eq. (14)]. The second type comes from the contributions of the direct interactions of the momentum of the flux tube with the quarks, represented by the spacelike potentials $\mathbf{A}_1$ and $\mathbf{A}_2$ [Eqs. (16)]. The latter terms induce negative signs to the spin-orbit couplings, in opposite direction to the former one, a feature which is also observed on phenomenological grounds for the large-distance effects in fine splitting.

The relativistic corrections to the interquark potential arising from the Wilson loop were analyzed and evaluated in the literature by Eichten and Feinberg [21], Gromes [22], Brambilla, Prosperi *et al.* [18, 19, 20], Brambilla, Pineda, Soto and Vairo [23].

The Wilson loop approach was also used for the study of quarkonium systems by Dosch, Simonov *et al.* with the use of the stochastic vacuum model [24].

In conclusion, the saturation of the Wilson loop averages in the large-N_c limit by minimal surfaces provides a systematic tool for investigating the large-distance dynamics of quarkonium systems. A complete study of those systems necessitates the incorporation of the short-distance contributions, taken into account by perturbation theory.

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